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ROLE OF CLOUD RENORMALIZATION IN CONVOLUTION MODELS FOR EMC AND DRELL-YAN RATIOS

Michael Buballa

Department of Physics, State University of New York,

Stony Brook, New York 11794, U.S.A.

Abstract

Generalizing a recent convolution model of the nucleon which explicitly conserves charge and momentum we reexamine the influence of an enhanced pion field in nuclei on EMC and Drell-Yan ratios. Due to wave function renormalization constants the effect is more than 50% smaller than predicted by the standard pion excess model. In particular there is no discrepancy between the EMC data and our results in the x-region which was expected to be most sensitive to the pion.

When in 1983 the EMC collaboration found an enhanced nucleon structure function $F_2(x)$ in iron at $x \leq 0.3$ [1], this was interpreted as a signature for an enhanced pion field in nuclei which was predicted already in the 70s [2]. Starting from the pion distribution function of the free nucleon, calculated in 1974 by Sullivan [3], Llewellyn Smith [4] was able to explain the effect by the assumption that the iron nucleus contains about 0.15 additional pions per nucleon. Generalizing Sullivan's expression Ericson and Thomas [5] showed that these extra pions result in a rather natural way from the attractive pion-nucleon interaction in a conventional meson-exchange description of nuclear matter. Consequently, after the enhancement of F_2 could not be confirmed by later experiments [6, 7], it was concluded that there are no excess pions in nuclei, calling into question the meson-exchange picture itself [8]. These conclusions were corroborated by the non-existence of a nuclear dependence of Drell-Yan cross sections [9]. Recent attempts to solve the problem are based on a possible change of the gluon properties in nuclei [8] or the partial restoration of chiral symmetry with density [10]. In this communication we do not intend to argue in favor or against these suggestions. Instead we want to reinvestigate the question whether an enhanced pion field – if it exists – really leads to a detectable enhancement of the EMC and Drell-Yan ratios. In particular we want to discuss the role of wavefunction renormalization constants which have not been considered in the past.

Starting point of our investigations is the two-phase model for the nucleon developed by Szczurek and Speth [11]. In this model the nucleon consists of a bare core and a meson-baryon phase. Here we want to restrict ourselves to pion and nucleon degrees of freedom which are supposed to be the main players in the region $0.1 \leq x \leq 0.3$. Then the distribution function for a quark with flavor f in a proton reads:

$$q_f^p(x) \;=\; Z_N\{q_{f,bare}^p(x) \;+\; \sum_i c_i\,[\,\int_x^1 \frac{dy}{y}\; f^{N/N}(y)\; q_{f,bare}^{N_i}(\frac{x}{y}) \;+\; \int_x^1 \frac{dy}{y}\; f^{\pi/N}(y)\; q_f^{\pi_i}(\frac{x}{y})\,]\,\}\;. \tag{1}$$

For the sake of readability we do not write the Q^2-dependence of the distribution functions explicitly. The two convolution integrals are due to the pion-nucleon phase or, more precisely, a $\pi^+ n$ phase with $c_i = 2/3$ and a $\pi^o p$ phase with $c_i = 1/3$. The function $f^{\pi/N}$

is the Sullivan expression for the momentum distribution of the pion while $f^{N/N}$ is the analogous distribution function of the recoil nucleon. The sum is multiplied by an overall constant Z_N which normalizes the total probability of finding the dressed nucleon either in the bare nucleon phase or in the pion-nucleon phase to unity. It is given by

$$Z_N^{-1} - 1 \;=\; M_1^{\pi/N} \equiv \int_0^1 dy\; f^{\pi/N}(y) \;=\; \int_0^1 dy\; f^{N/N}(y)\;. \tag{2}$$

This is a consequence of the well-known number sum rules for valence quarks reflecting the conservation of baryon number, charge and strangeness:

$$\int_0^1 dx\;(q_f^p(x) - \bar{q}_f^p(x)) \;=\; \int_0^1 dx\;(q_{f,bare}^p(x) - \bar{q}_{f,bare}^p(x)) \;=\; N_f^p\;, \tag{3}$$

with $N_u^p = 2$, $N_d^p = 1$ and $N_s^p = 0$. Note that eq. (2) also implies that the first moments of the distribution functions $f^{\pi/N}$ and $f^{N/N}$ have to be equal.

Similarly momentum conservation

$$\sum_f \int_0^1 dx\; x\, q_f^p(x) \;=\; \sum_f \int_0^1 dx\; x\, q_{f,bare}^p(x) \tag{4}$$

leads to a relation between the second moments of $f^{\pi/N}$ and $f^{N/N}$ [12]. In eq. (4) the flavor index f runs over quarks and antiquarks. In principle gluons have to be included, too. However, since experimentally the momentum fraction carried by quarks is about the same in a pion as it is in a nucleon ($\sim 47\%$ at $Q^2 = 50(GeV/c)^2$) we must demand momentum conservation for quarks and gluons separately when we want to keep the internal structure of the constituents unchanged.

Higher moments of $f^{N/N}$ are not related to $f^{\pi/N}$ by conservation laws. However, according to the intuitive picture that the nucleon breaks up into a pion which carries the fraction y of the plus momentum of the original nucleon and a recoil nucleon which carries the missing fraction $1-y$ it is frequently assumed [11]

$$f^{N/N}(y) \;=\; f^{\pi/N}(1-y). \tag{5}$$

This automatically guarantees the validity of eqs. (2) - (4) but is not a necessary consequence of these relations.

Many features we want to discuss later on can be expressed in terms of the mean number of pions per nucleon, $\langle n_\pi \rangle$. For the free nucleon we get from eqs. (1) and (2):

$$\langle n_\pi^N \rangle \;=\; Z_N\, M_1^{\pi/N} \;=\; \frac{M_1^{\pi/N}}{1\,+\,M_1^{\pi/N}}\;. \tag{6}$$

Note that even for large moments $\langle n_\pi \rangle$ is always less than one, in agreement with the two-phase picture of the nucleon we have started with: Since at any time there is at most one pion, $\langle n_\pi \rangle$ should never exceed unity. (Of course, for $M_1^\pi > 1$, i.e. when the $N\pi$ loop becomes more important than the bare nucleon phase the two-phase model itself becomes unrealistic and multiple pion phases have to be included. In practice, however, this does not happen.)

The essential point of eq. (1) is that the renormalization constant Z_N is an *overall* prefactor, i.e. not only the bare nucleon part but also the cloud contribution becomes renormalized. In contrast Melnitchouk and Thomas [12] argue that one should multiply only the bare nucleon contribution by a normalization constant which is given by $\tilde{Z}_N = 1 - M_1^{\pi/N}$. This is of course what one gets when eq. (1) is expanded into a power series in the πNN coupling constant until order $g_{\pi NN}^2$. Numerically however, this makes a big difference. Note that in this case the mean number of pions per nucleon, $\langle n_\pi \rangle$ is identical with the first moment M_1^π. The difference between ref. [12] and ref. [11] (i.e. our eq. (1)) is due to a different summation scheme of diagrams when the dressed nucleon propagator is calculated. The authors of ref. [12] perform a systematic expansion in the coupling constant until quadratic order, thus taking into account only single one-loop diagrams. In contrast, the authors of ref. [11] iterate the meson-baryon loops until infinite order. Very recently Holtmann showed by a comparison of the corresponding next order corrections that this is a significantly better approximation [13]. In any case, when we turn to the calculation of nuclear structure functions a strict expansion in the πNN coupling constant is no longer a useful scheme. The medium modifications of the pion distribution function are mainly due to rescattering processes and Pauli blocking. So if we restrict ourselves to diagrams of order $g_{\pi NN}^2$ we should also drop all rescattering diagrams. In this case we would be left with Pauli blocking as the only medium effect which would lead to a

depletion rather than an enhancement of the pion field. On the other hand, if we sum over an infinite set of rescattering diagrams as it is usually done [5, 16] we should also iterate the πN loops which dress the nucleon. This means we should normalize like in eq. (1).

In the pion excess model [5, 16] it is assumed that the quark distribution functions of the bare nucleon and of the pion remain unchanged inside a nucleus. Thus in analogy to eq. (1) we write

$$q_f^{p/A}(x) \;=\; Z_A\{q^p_{f,bare}(x) \;+\; \sum_i c_i\; [\; \int_x^A \frac{dy}{y}\; f^{N/A}(y)\; q^{N_i}_{f,bare}(\frac{x}{y}) \;+\; \int_x^A \frac{dy}{y}\; f^{\pi/A}(y)\; q_f^{\pi_i}(\frac{x}{y})\;]\;\}, \tag{7}$$

with relations analogous to eqs. (2) - (4) in order to guarantee baryon number, charge and momentum conservation. Note that we still define the Bjorken variable x with respect to the nucleon mass. Therefore, unlike the free nucleon case, $f^{\pi/A}$ and $f^{N/A}$ do not need to vanish for $y \geq 1$ and in general a relation like eq. (5) can not be correct in a nucleus.

On the other hand in microscopic models the upper limit y_{max} for non-vanishing $f^{\pi/A}$ does not really become equal to the nuclear mass number A as we have indicated in eq. (7) but remains rather close to unity. For the model presented in ref. [10] for instance, we find $y_{max} = 1 + \frac{k_F^2}{m_N^2}$ which is about 1.1 at nuclear matter density (The model of refs. [5, 16] even leads to $y_{max} < 1$ but that is an artefact of the nonrelativistic kinematics.). Furthermore the shape of the pion distribution function is not dramatically changed by nuclear effects and the main contribution still comes from the regime $y \leq 1$. Thus before we turn to a more realistic model we consider a schematic model where the pion distribution function in a nucleus is simply given by $f^{\pi/N}$ multiplied by an enhancement factor λ:

$$f^{\pi/A}(y) \;=\; \lambda\, f^{\pi/N}(y)\;. \tag{8}$$

For the distribution function of the recoil nucleons $f^{N/A}(y)$ we assume that a relation analogous to eq. (5) holds.

Within this model we calculate EMC and Drell-Yan ratios for different values of λ. For the πNN formfactor we choose a monopole parameterization with a cutoff parameter

$\Lambda_{\pi NN} = 800 MeV$. We assume the sea quark distributions of the bare nucleons and of the pions to be $SU(3)$-symmetric and the proton and the neutron as well as the different pions to be related by isospin rotations. We parameterize the quark distributions of the bare proton in the following simple form:

$$x\, q^p_{f,bare}(x) \;=\; N_f\, x^{\alpha_f}\, (1-x)^{\beta_f} \quad . \tag{9}$$

The parameters are chosen such that the experimental distribution functions of the dressed proton [14] are roughly reproduced by eq. (1). We find $\alpha_f = 0.65$ and $\beta_f = 3$ for valence up quarks, $\alpha_f = 0.65$ and $\beta_f = 4$ for valence down quarks and $\alpha_f = 0$, $\beta_f = 9$ and $N_f = 0.17$ for sea quarks. The normalization of the valence quarks is of course fixed by eq. (3). For the pion we take the experimental quark distribution functions of ref. [15].

The left panel of fig. 1 shows our results for the EMC ratio F_2^A/F_2^D for isospin symmetric nuclear matter. For the deuteron nuclear effects are neglected. At small values of x the ratio becomes enhanced with increasing λ whereas it becomes reduced at larger x. At $x \simeq 0.2$ the EMC ratio remains unity independently of λ. This behavior is easy to understand. Using $F_2(x) = x \sum e_f^2 q_f(x)$ we find for the EMC ratio at $x = 0$:

$$\frac{F_2^A(0)}{F_2^N(0)} \;=\; 1 \;+\; (\langle n_\pi^A\rangle - \langle n_\pi^N\rangle)\, \frac{F_2^\pi(0)}{F_2^N(0)}\, . \tag{10}$$

This result is quite general and does not depend on the specific assumption eq. (8). For our schematic model we can also calculate the ratio at $x = 1$. At this point the contributions of the pion and the recoil nucleon vanish and the EMC ratio is just the ratio of the renormalization constants:

$$\frac{F_2^A(1)}{F_2^N(1)} \;=\; \frac{Z_A}{Z_N} \;=\; \frac{1-\langle n_\pi^A\rangle}{1-\langle n_\pi^N\rangle}\, . \tag{11}$$

Thus with an enhanced pion field ($\langle n_\pi^A\rangle > \langle n_\pi^N\rangle$) the EMC ratio is greater than one at $x = 0$ and less than one at $x = 1$. So somewhere in between there must be a point x_o with $F_2^A(x_o) = F_2^N(x_o)$. In our schematic model this point is given by by the condition

$$F_2^N(x_o) \;=\; F^N_{2,bare}(x_o) \tag{12}$$

independently of λ. We find $x_o \sim 0.2$. That means an enhanced pion field leads to an EMC ratio greater than one only for $x \leq 0.2$. In order to estimate whether this is a detectable effect we also plotted experimental data points in fig. 1. Note that at $x \leq 0.1$ nuclear shadowing sets in which is not included in our model. So there is only a very small window $0.1 \leq x \leq 0.2$ where an enhancement could be seen. For realistic parameters ($\lambda \leq 2$) this seems to be very difficult.

Here the wave function renormalization plays an important role. As we can see from eq. (10) the enhancement is directly related to the number of excess pions which is given by

$$\langle n_\pi^A \rangle \; - \; \langle n_\pi^N \rangle \; = \; Z_N Z_A \, (M_1^{\pi/A} - M_1^{\pi/N}). \tag{13}$$

Thus compared with calculations without normalization factors the effect is reduced by a factor $Z_N Z_A$. Since $Z_A < Z_N$ for an enhanced pion field and $Z_N \simeq 0.7$ for our parameters, the additional pion contribution is less than half of what it would be without normalization factors. It is also crucial that the renormalization includes the pion cloud, like in eq. (1). If we renormalized only the bare nucleon, like in ref. [12], we still would get eqs. (10) - (12). In this case, however, the pion numbers $\langle n_\pi \rangle$ would be identical with the first moments M_1^π and we would find eq. (13) without the factor $Z_A Z_N$.

In the right panel of fig. 1 we show the corresponding Drell-Yan ratios as a function of x_t, the Bjorken variable for the target nucleus. The Bjorken variable for the projectile was kept fixed at $x_p = 0.5$. At this large value of x_p we are mainly sensitive to the target sea which is dominated by the pion contribution. Therefore this process is much more sensitive to detect an enhancement of the pion field although the effect is also strongly reduced by the normalization factors.

After these schematic studies we want to make the model somewhat more realistic. The enhancement of the pion distribution function in nuclear matter has been calculated many times [5, 16, 10]. In these models the bare pion propagator becomes dressed by nucleon hole and Δ hole excitations. As we have mentioned before, $f^{\pi/A}(y)$ does not have to vanish for $y \geq 1$ since a pion may carry more than the fraction $1/A$ of the plus

momentum of the nucleus. Therefore conservation of charge and momentum cannot be enforced by a relation like eq. (5). This can also be seen from more physical arguments: Since we are summing an infinite series of particle hole polarizations we should allow the virtual photon not only to couple to the recoil nucleon which corresponds to the first emitted pion but also to particle states which are created by absorbing a pion, i.e. to nucleons which have interacted with another nucleon before. The most natural way to describe all these processes effectively is to incorporate them into Fermi motion and nuclear binding.

We proceed as follows: Integrating eqs. (1) and (7) over x and using charge conservation, i.e. eq. (2) and the corresponding relation for the nuclear case, we can eliminate the recoil nucleon distributions as well as the distribution functions of the bare core. In this way we find expressions for $\int dx\, q_f^{p/A}(x)$ which only depend on the distribution functions of the (dressed) free nucleon and the pion distribution functions $f^{\pi/N}$ and $f^{\pi/A}$. By construction the integrands already satisfy the requirements of charge conservation, but momentum is not yet conserved. In the next step we convolute the nucleonic part with a distribution function $f^N_{Fermi}(z)$ which describes the effects of Fermi motion and nuclear binding. This procedure does not change the first moments of the quark distribution functions but it changes the second moments as a function of a binding parameter η. Thus by tuning η to the appropriate value we can achieve momentum conservation while keeping the charge to be conserved.

We find the following expressions for the quark distribution functions of a nuclear proton:

$$q_f^{p/A}(x) = \int_x^A \frac{dz}{z}\, f^N_{Fermi}(z)\, q_f^p(\frac{x}{z}) + \sum_i c_i \int_x^A \frac{dy}{y}\, (Z_A\, f^{\pi/A}(y) - Z_N\, f^{\pi/N}(y))\, q_f^{\pi_i}(\frac{x}{y}) - \Delta_f^p(x), \tag{14}$$

with

$$\Delta_u^p(x) = \frac{2}{3}\, (Z_N - Z_A)[\int_x^A \frac{dz}{z}\, f^N_{Fermi}(z)(u^p(\frac{x}{z}) - u^n(\frac{x}{z}))$$

$$+\frac{4}{3}Z_N\int_x^1\frac{dy}{y}\,f^{\pi/N}(y)(u^{\pi^+}(\frac{x}{y})-u^{\pi^o}(\frac{x}{y}))]\,, \tag{15}$$

$\Delta_d^p(x)=-\Delta_u^p(x)$ and $\Delta_f^p(x)=0$ in all other cases.

Except for the normalization factors the two integrals in eq. (14) are identical with the standard pion excess model, as discussed e.g. in ref. [16]. While the first integral describes the redistribution of the nucleons keeping the baryon number fixed, the second integral is the contribution from the excess pions. However, since the pion cloud around a proton carries a positive net charge, in order to conserve the total charge we have to correct for the extra charge of the excess pions: Due to the emission of additional π^+ mesons the number of up quarks associated with the baryonic part of a nuclear proton is lowered while the number of down quarks is enhanced. This is the origin of Δ_f^p in eq. (14). Since this term does not contribute to the momentum balance, only its first moment is fixed and one can easily find other expressions which do the job as well as eq. (15). In practice, however, this does not really matter: Since for neutrons everything is just the other way around, in isospin symmetric nuclear matter the Δ_f-terms completely drop out. In nuclei with a small neutron excess they persist, but they are strongly suppressed. In fact, for ^{56}Fe we find that the influence of Δ_f can be almost neglected. Thus the main difference to earlier calculations [5, 16] is the appearance of normalization factors in eq. (14) which reduce the pion enhancement in the same way as we have discussed for the schematic model.

Our results for the EMC and Drell-Yan ratios (^{56}Fe compared with deuterium) are shown in fig. 2. The solid lines represent the ratios obtained from eq. (14) imposing momentum conservation. We take the experimental quark distributions of Owens for the pion [15] as well as for the nucleon [14]. The nuclear pion distribution function is calculated microscopically within the framework of a nuclear matter RPA calculation, including NN^{-1} as well as ΔN^{-1} excitations. Short-range correlations are taken into account by folding the potential with a correlation function $g_c(r)=1-j_o(q_cr)$. In the NN channel we choose $q_c=3.94fm^{-1}$ corresponding to the mass of the ω-meson. This

gives rise to a momentum dependent Migdal parameter with the limit $g'_{NN}(q=0) = 0.57$. In the $N\Delta$ and the $\Delta\Delta$ channel q_c is tuned to reproduce $g'_{N\Delta}(0) = g'_{\Delta\Delta}(0) = 1/3$, the classical Lorentz-Lorenz value. For details of the interaction see ref. [10], although in the present work we do not introduce density dependent masses and coupling constants which are discussed there. We use the nucleon distribution function $f^N_{Fermi}(z)$ derived from a Fermi gas model using the correct relativistic normalization [17]. For the Fermi momentum we take $k_F = 260 MeV$ which corresponds to $\rho = 0.87\rho_o$, the average density of the ^{56}Fe nucleus. The parameter η is fixed by momentum conservation. In order to compensate for the additional momentum carried by the pions ($\sim 4.5\%$) we have to choose $\eta = 0.94$.

For comparison we also show the result obtained within the schematic model we discussed above (dotted line). We choose $\lambda = 1.75$ which corresponds to the enhancement of the pion field we find in the microscopic calculation. Of course, the schematic model is too simplistic to produce the rise of the EMC ratio seen at $x \geq 0.7$. However, for $x \leq 0.6$ and in particular for $x \leq 0.3$ – our region of interest – the two models agree quite well. This is also true more or less for the Drell-Yan ratio. Moreover, in a more detailed investigation (not shown in fig. 2) we find that a major part of the differences comes about from the fact that we do not exactly start from the same quark distribution functions of the free nucleon. Remember, in the schematic model we calculated $q^p_f(x)$ from eq. (1) using the simple parameterization eq. (9), whereas we directly used the experimental parameterizations [14] in the more realistic model. This is also the reason why the ratios do not exactly agree at $x = 0$ where they are fixed by eq. (10). So with some more effort in finding a good parameterization for the quark distribution functions of the bare nucleon the dotted line would come even closer to the solid curve.

The comparison shows that for this type of models the EMC ratio at small and intermediate x is to a very large extent determined by the number of excess pions and the amount of momentum transferred from the nucleons to the pions: The first fixes the ratio at $x = 0$ while the latter determines the slope of the drop with increasing x. Details

which are related to higher moments of the nucleon and the pion distribution functions are less important. Therefore we believe that the major conclusions we have drawn from the schematic model remain valid for more sophisticated descriptions. In particular it will be very difficult to produce any significant enhancement in the region $0.1 \leq 0.3$.

As mentioned above, in order to conserve the total momentum we have to choose the binding parameter to be $\eta = 0.94$. This corresponds to an average separation energy of $(1 - \eta)m_N \simeq 56MeV$ which is of course unrealistically large. In order to check the sensitivity to this parameter we give up the requirement of momentum conservation and choose $\eta = 0.97$, the value proposed by Li et al. [17]. The results are shown as the dashed dotted curves in fig. 2. Since we have not changed the number of excess pions the ratios at $x = 0$ remain unchanged, but the EMC ratio drops off more slowly. However, for $x \leq 0.4$, i.e. in the region where the pion enhancement was supposed to show up, there is still no disagreement with the data.

The momentum conserving EMC result (solid line) looks very similar to fig. 4 of ref. [16]. Note however, that our short-range repulsion is much weaker and more realistic [10]. Without wave function renormalization this would lead to a much larger number of excess pions (as discussed in eq. (13)) and consequently to a larger ratio at small x. In order to demonstrate this we replace Z_N and Z_A in eq. (14) by unity. The results are indicated by the long and the short-dashed lines in fig. 2. When we enforce momentum conservation by choosing $\eta = 0.88$ we get the results indicated by the short-dashed line. Similar to the behavior discussed in fig. 1 the EMC ratio drops off very fast and becomes unity at $x \sim 0.2$. When we choose the realistic parameter $\eta = 0.97$ we arrive at the long-dashed curve. This is the only case where the EMC-data are significantly overestimated in the whole area below $x \leq 0.3$.

To conclude, we find that EMC experiments are not sufficiently sensitive to detect an effect due to a possibly enhanced pion field in nuclei. Within a schematic model we show that the enhancement of the pion distribution function must be unrealistically large to produce a significant deviation from the data in the region $0.1 \leq x \leq 0.3$. These findings

are confirmed by more realistic calculations. Our results are based on a generalization of the two-phase model for the free nucleon of Szczurek and Speth [11]. In this model the pion contribution is strongly reduced by a renormalization constant. Other prescriptions which renormalize the bare nucleon part only are based on a power expansion in the coupling constant and seem to be unsuitable for nuclear matter.

In contrast to EMC experiments, Drell-Yan experiments are sensitive to conclude that there is indeed no enhancement of the pion field in nuclei. Although the renormalization constants reduce the discrepancy by more than 50% the data are still overestimated in a standard RPA calculation. This leaves room for more unconventional explanations [10].

So far we have restricted ourselves to nucleon and pion degrees of freedom. Although heavier mesons are known to be more important in the intermediate regime $x \sim 0.5$, through the normalization factors and the momentum balance they also have some influence on the lower x region. We do not expect, however, that this will spoil our main conclusions.

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Figure Captions

Fig. 1 EMC ratio (left panel) and Drell-Yan ratio (right panel) for isospin symmetric nuclear matter calculated within the schematic model (eq. (8)) with $\lambda = 1$ (solid line), $\lambda = 1.5$ (dashed-dotted line), $\lambda = 2$ (dashed line) and $\lambda = 4$ (dotted line). The data points [6, 7, 9] are plotted in order to give an estimate for the experimental sensitivity.

Fig. 2 EMC ratio (left panel) and Drell-Yan ratio (right panel) for ^{56}Fe calculated in a realistic model using eq. (14). The solid line shows the result for the binding parameter $\eta = 0.94$, obtained from momentum conservation, while the dashed-dotted line is calculated with the more realistic value $\eta = 0.97$. The dashed lines indicate the results we get in a calculation without renormalizing the pion contribution ($Z_A = Z_N = 1$), short-dashed line: $\eta = 0.88$, long-dashed line: $\eta = 0.97$. For comparison we also show the result obtained within the schematic model with $\lambda = 1.75$ (dotted line). The EMC data were taken from refs. [6] (^{40}Ca) and [7] (^{56}Fe), the Drell-Yan data were taken from ref. [9] (^{56}Fe).